

# Low scale thermal leptogenesis in neutrinophilic Higgs doublet models

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## Abstract

It is well-known that leptogenesis in low energy scale is difficult in the conventional Type-I seesaw mechanism with hierarchical right-handed neutrino masses. We show that in a class of two Higgs doublet model, where one Higgs doublet generates masses of quarks and charged leptons whereas the other Higgs doublet with a tiny vacuum expectation value generates neutrino Dirac masses, large Yukawa couplings lead to a large enough CP asymmetry of the right-handed neutrino decay. Thermal leptogenesis suitably works at low energy scale as keeping no enhancement of lepton number violating wash out effects. We will also point out that thermal leptogenesis works well without confronting gravitino problem in a supersymmetric neutrinophilic Higgs doublet model with gravity mediated supersymmetry breaking. Neutralino dark matter and baryon asymmetry generation by thermal leptogenesis are easily compatible in our setup.

## §1. Introduction

An origin of cosmological baryon asymmetry is one of the prime open questions in particle physics as well as in cosmology. Among various mechanisms of baryogenesis, leptogenesis<sup>1)</sup> is one of the most attractive idea because of its simplicity and the connection to neutrino physics. Particularly, thermal leptogenesis requires only the thermal excitation of heavy right-handed Majorana neutrinos which generate tiny neutrino masses via the seesaw mechanism<sup>2)</sup> and provides several implications for the light neutrino mass spectrum.<sup>3)</sup> The size of CP asymmetry in a right-handed neutrino decay is, roughly speaking, proportional to the mass of right-handed neutrino. Thus, we obtain only insufficiently small CP violation for a lighter right-handed neutrino mass. That is the reason why it has been regarded that leptogenesis in low energy scale is in general difficult in the conventional Type I seesaw mechanism.<sup>4), 5)</sup>

On the other hand, in supersymmetric models with conserved R-parity to avoid rapid proton decay, thermal leptogenesis faces with “gravitino problem” that the overproduction of gravitinos spoils the success of Big Bang Nucleosynthesis (BBN),<sup>6)</sup> whereas the stable lightest supersymmetric particle (LSP) becomes dark matter candidate. In order not to overproduce gravitinos, the reheating temperature after inflation should not be so high that right-handed neutrinos can be thermally produced.<sup>7)</sup> In the framework of gravity mediated supersymmetry (SUSY) breaking, a few solutions, e.g., gravitino LSP with R-parity violation,<sup>8)</sup> very light axino LSP<sup>9)</sup> and strongly degenerated right-handed neutrino masses,<sup>10)</sup> have been proposed.

Recently, a new class of two Higgs doublet models (THDM)<sup>11)</sup> has been considered in Refs.<sup>12)–17)</sup> The motivation is as follows. As mentioned above, seesaw mechanism naturally realizes tiny masses of active neutrinos through heavy particles coupled with left-handed neutrinos. However, those heavy particles are almost decoupled in the low-energy effective theory, few observations are expected in collider experiments. Then, some people consider a possibility of reduction of seesaw scale to TeV,<sup>18), 19)</sup> where effects of TeV scale right-handed neutrinos might be observed in collider experiments such as Large Hadron Collider (LHC) and International Linear Collider (ILC). However, they must introduce a fine-tuning in order to obtain both tiny neutrino mass and detectable left-right neutrino mixing through which experimental evidences can be discovered. Other right-handed neutrino production processes in extended models by e.g.,  $Z'$  exchange<sup>20)</sup> or Higgs/Higgsino decay<sup>21)</sup> also have been pointed out. Here, let us remind that Dirac masses of fermions are proportional to their Yukawa couplings as well as a vacuum expectation value (VEV) of the relevant Higgs field. Hence, the smallness of a mass might be due to not a small Yukawa coupling but a small VEV of the Higgs field. Such a situation is indeed realized in some THDM. For

example, in Type-II THDM with a large  $\tan\beta$ , the mass hierarchy between up-type quark and down-type quark can be explained by the ratio of Higgs VEVs, and when  $\tan\beta \sim 40$ , Yukawa couplings of top and bottom quark are same scale of order of unity.<sup>22)</sup> Similarly, there is a possibility that smallness of the neutrino masses comparing to those of quarks and charged leptons is originating from an extra Higgs doublet with the tiny VEV. This idea is that neutrino masses are much smaller than other fermions because the origin of them comes from different VEV of different Higgs doublet, and then we do not need extremely tiny neutrino Yukawa couplings. Let us call this kind of model<sup>12)–17)</sup> *neutrinophilic Higgs doublet model*. Especially, in models in Refs.,<sup>12), 17)</sup> tiny Majorana neutrino masses are obtained through a TeV scale Type-I seesaw mechanism without requiring tiny Yukawa couplings.

Notice that neutrino Yukawa couplings in neutrinophilic Higgs doublet models do not need to be so small. This fact has significant implication to leptogenesis. The CP violation of right-handed neutrino decay is proportional to neutrino Yukawa coupling squared. We can obtain a large CP violation for such large neutrino Yukawa couplings. This opens new possibility of low scale thermal leptogenesis. In this paper, we will show that CP asymmetry is enhanced and thermal leptogenesis suitably works in multi-Higgs models with a neutrinophilic Higgs doublet field, where the tiny VEV of the neutrinophilic Higgs field has equivalently larger neutrino Yukawa couplings, and then TeV-scale seesaw works well. We will show that the thermal leptogenesis suitably works at low energy scale as avoiding enhancement of lepton number violating wash out effects. We will also point out that thermal leptogenesis in gravity mediated SUSY breaking works well without confronting gravitino problem in a supersymmetric model.

## §2. Neutrinophilic Higgs doublet models

### 2.1. Minimal neutrinophilic THDM

Let us show a two Higgs doublet model, which we call neutrinophilic THDM model, originally suggested in Ref.<sup>12)</sup> In the model, one additional Higgs doublet  $\Phi_\nu$ , which gives only neutrino Dirac masses, besides the SM Higgs doublet  $\Phi$  and a discrete  $Z_2$ -parity are introduced. The  $Z_2$ -symmetry charges (and also lepton number) are assigned as the following table. Under the discrete symmetry, Yukawa interactions are given by

$$\mathcal{L}_{\text{Yukawa}} = y^u \bar{Q}_L \Phi U_R + y^d \bar{Q}_L \tilde{\Phi} D_R + y^l \bar{L} \Phi E_R + y^\nu \bar{L} \Phi_\nu N + \frac{1}{2} M \bar{N}^c N + \text{h.c.} \quad (2.1)$$

where  $\tilde{\Phi} = i\sigma_2 \Phi^*$ , and we omit a generation index.  $\Phi_\nu$  only couples with  $N$  by the  $Z_2$ -parity so that flavor changing neutral currents (FCNCs) are suppressed. Quark and charged lepton sectors are the same as Type-I THDM, but notice that this neutrinophilic THDM is quite

fields	$Z_2$ -parity	lepton number
SM Higgs doublet, $\Phi$	+	0
new Higgs doublet, $\Phi_\nu$	−	0
right-handed neutrinos, $N$	−	1
others	+	$\pm 1$ : leptons, 0: quarks

different from conventional Type-I, II, X, Y THDMs.<sup>11)</sup>

The Higgs potential of the neutrinophilic THDM is given by

$$\begin{aligned}
V^{\text{THDM}} = & m_\Phi^2 \Phi^\dagger \Phi + m_{\Phi_\nu}^2 \Phi_\nu^\dagger \Phi_\nu - m_3^2 (\Phi^\dagger \Phi_\nu + \Phi_\nu^\dagger \Phi) + \frac{\lambda_1}{2} (\Phi^\dagger \Phi)^2 + \frac{\lambda_2}{2} (\Phi_\nu^\dagger \Phi_\nu)^2 \\
& + \lambda_3 (\Phi^\dagger \Phi) (\Phi_\nu^\dagger \Phi_\nu) + \lambda_4 (\Phi^\dagger \Phi_\nu) (\Phi_\nu^\dagger \Phi) + \frac{\lambda_5}{2} [(\Phi^\dagger \Phi_\nu)^2 + (\Phi_\nu^\dagger \Phi)^2]. \quad (2.2)
\end{aligned}$$

The  $Z_2$ -symmetry is softly broken by  $m_3^2$ . Taking a parameter set,

$$m_\Phi^2 < 0, \quad m_{\Phi_\nu}^2 > 0, \quad |m_3^2| \ll m_{\Phi_\nu}^2, \quad (2.3)$$

we can obtain the VEV hierarchy of Higgs doublets,

$$v^2 \simeq \frac{-m_\Phi^2}{\lambda_1}, \quad v_\nu \simeq \frac{-m_3^2 v}{m_{\Phi_\nu}^2 + (\lambda_3 + \lambda_4 + \lambda_5)v^2}, \quad (2.4)$$

where we have decomposed the SM Higgs doublet  $\Phi$  and the extra Higgs doublet  $\Phi_\nu$  as

$$\Phi = \begin{pmatrix} v + \frac{1}{\sqrt{2}}\phi^0 \\ \phi^- \end{pmatrix}, \quad \Phi_\nu = \begin{pmatrix} v_\nu + \frac{1}{\sqrt{2}}\phi_\nu^0 \\ \phi_\nu^- \end{pmatrix}. \quad (2.5)$$

When we take values of parameters as  $m_\Phi \sim 100$  GeV,  $m_{\Phi_\nu} \sim 1$  TeV, and  $|m_3^2| \sim 10$  GeV<sup>2</sup>, we can obtain  $v_\nu \sim 1$  MeV. The smallness of  $|m_3^2|$  is guaranteed by the “softly-broken”  $Z_2$ -symmetry.

For a very large  $\tan \beta = v/v_\nu (\gg 1)$  limit we are interested in, the five physical Higgs boson states and those masses are respectively given by

$$H^\pm \simeq [\phi_\nu^\pm], \quad m_{H^\pm}^2 \simeq m_\nu^2 + \lambda_3 v^2, \quad (2.6)$$

$$A \simeq \text{Im}[\phi_\nu^0], \quad m_A^2 \simeq m_\nu^2 + (\lambda_3 + \lambda_4 + \lambda_5)v^2, \quad (2.7)$$

$$h \simeq \text{Re}[\phi^0], \quad m_h^2 \simeq 2\lambda_1 v^2, \quad (2.8)$$

$$H \simeq \text{Re}[\phi_\nu^0], \quad m_H^2 \simeq m_\nu^2 + (\lambda_3 + \lambda_4 + \lambda_5)v^2, \quad (2.9)$$

where negligible  $\mathcal{O}(v_\nu^2)$  and  $\mathcal{O}(m_3^2)$  corrections are omitted. Notice that  $\tan \beta$  is extremely large so that the SM-like Higgs  $h$  is almost originated from  $\Phi$ , while other physical Higgs

particles,  $H^\pm, H, A$ , are almost originated from  $\Phi_\nu$ . Since  $\Phi_\nu$  has Yukawa couplings only with neutrinos and lepton doublets, remarkable phenomenology can be expected which is not observed in other THDMs. For example, lepton flavor violation (LFV) processes and oblique corrections are estimated in Ref.<sup>(12)</sup> and charged Higgs processes in collider experiments are discussed in Refs.<sup>(15), (16)</sup> \*).

The neutrino masses including one-loop radiative corrections<sup>(14)</sup> are estimated as

$$(m_\nu)_{ij} = \sum_k \frac{y_{ik}^\nu v_\nu y_{kj}^{\nu T}}{M_k} + \sum_k \frac{y_{ik}^\nu y_{kj}^{\nu T} M_k}{16\pi^2} \left[ \frac{m_R^2}{m_R^2 - M_k^2} \ln \frac{m_R^2}{M_k^2} - \frac{m_I^2}{m_I^2 - M_k^2} \ln \frac{m_I^2}{M_k^2} \right], \quad (2.10)$$

where  $m_R$  and  $m_I$  are the masses of  $\text{Re}[\phi^0]$  and  $\text{Im}[\phi_\nu^0]$  respectively. It is easy to see the tree level contribution gives  $\mathcal{O}(0.1)$  eV neutrino masses for  $M_k \sim 1$  TeV,  $v_\nu \sim 1$  MeV and  $y^\nu = \mathcal{O}(1)$ . The one-loop contribution is induced for a nonvanishing  $\lambda_5$ . When  $m_R^2 - m_I^2 = 2\lambda_5 v^2 \ll m_0^2 = (m_R^2 + m_I^2)/2$ ,

$$(m_\nu)_{ij} = \frac{\lambda_5 v^2}{8\pi^2} \sum_k \frac{y_{ik}^\nu y_{jk}^\nu M_k}{m_0^2 - M_k^2} \left[ 1 - \frac{M_k^2}{m_0^2 - M_k^2} \ln \frac{m_0^2}{M_k^2} \right], \quad (2.11)$$

and it shows

$$(m_\nu)_{ij} = \frac{\lambda_5 v^2}{8\pi^2} \sum_k \frac{y_{ik}^\nu y_{jk}^\nu M_k}{M_k} \left[ \ln \frac{M_k^2}{m_0^2} - 1 \right], \quad (M_k^2 \gg m_0^2), \quad (2.12)$$

$$(m_\nu)_{ij} = \frac{\lambda_5 v^2}{8\pi^2 m_0^2} \sum_k y_{ik}^\nu y_{jk}^\nu M_k, \quad (m_0^2 \gg M_k^2), \quad (2.13)$$

$$(m_\nu)_{ij} \simeq \frac{\lambda_5 v^2}{16\pi^2} \sum_k \frac{y_{ik}^\nu y_{jk}^\nu}{M_k}, \quad (m_0^2 \simeq M_k^2). \quad (2.14)$$

Thus, when the masses of Higgs bosons (except for  $h$ ) and right-handed neutrinos are  $\mathcal{O}(1)$  TeV, light neutrino mass scale of order  $\mathcal{O}(0.1)$  eV is induced with  $\lambda_5 \sim 10^{-4}$ . Thus, whether tree-level effect is larger than loop-level effect or not is determined by the magnitude of  $\lambda_5$  (and  $m_A, m_H$ ), which contribute one-loop diagram.

## 2.2. A UV theory of neutrinophilic THDM

Here let us show a model in Ref.<sup>(17)</sup> as an example of UV theory of the neutrinophilic THDM. This model is constructed by introducing one gauge singlet scalar field  $S$ , which has a lepton number, and  $Z_3$ -symmetry shown as the following table. Under the discrete symmetry, Yukawa interactions are given by

$$\mathcal{L}_{\text{Yukawa}} = y^u \bar{Q}^L \Phi U_R + y^d \bar{Q}^L \tilde{\Phi} D_R + y^l \bar{L} \Phi E_R + y^\nu \bar{L} \Phi_\nu N + \frac{1}{2} y^N S \bar{N}^c N + \text{h.c.} \quad (2.15)$$

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\*) The model deals with Dirac neutrino version in neutrinophilic THDM, but phenomenology of charged lepton has a similar region in part.

fields	$Z_3$ -charge	lepton number
SM Higgs doublet, $\Phi$	1	0
new Higgs doublet, $\Phi_\nu$	$\omega^2$	0
new Higgs singlet, $S$	$\omega$	-2
right-handed neutrinos, $N$	$\omega$	1
others	1	$\pm 1$ : leptons, 0: quarks

The Higgs potential can be written as

$$\begin{aligned}
V = & m_\Phi^2 |\Phi|^2 + m_{\Phi_\nu}^2 |\Phi_\nu|^2 - m_S^2 |S|^2 - \lambda S^3 - \kappa S \Phi^\dagger \Phi_\nu \\
& + \frac{\lambda_1}{2} |\Phi|^4 + \frac{\lambda_2}{2} |\Phi_\nu|^4 + \lambda_3 |\Phi|^2 |\Phi_\nu|^2 + \lambda_4 |\Phi^\dagger \Phi_\nu|^2 \\
& + \lambda_S |S|^4 + \lambda_\Phi |S|^2 |\Phi|^2 + \lambda_{\Phi_\nu} |S|^2 |\Phi_\nu|^2 + h.c..
\end{aligned} \tag{2.16}$$

$Z_3$ -symmetry forbids dimension four operators,  $(\Phi^\dagger \Phi_\nu)^2$ ,  $\Phi^\dagger \Phi_\nu |\Phi|^2$ ,  $\Phi^\dagger \Phi_\nu |\Phi_\nu|^2$ ,  $S^4$ ,  $S^2 |S|^2$ ,  $S^2 |\Phi|^2$ ,  $S^2 |\Phi_\nu|^2$ , and dimension two or three operators,  $\Phi^\dagger \Phi_\nu$ ,  $S |\Phi|^2$ ,  $S |\Phi_\nu|^2$ . Although there might be introduced small soft breaking terms such as  $m_3^2 \Phi^\dagger \Phi_\nu$  to avoid domain wall problem, we omit them here, for simplicity. It has been shown that, with  $\kappa \sim 1$  MeV, the desirable hierarchy of VEVs

$$v_s \equiv \langle S \rangle \sim 1 \text{ TeV}, \quad v \sim 100 \text{ GeV}, \quad v_\nu \sim 1 \text{ MeV}, \tag{2.17}$$

and neutrino mass

$$m_\nu \simeq \frac{y^{\nu^2} v_\nu^2}{M_N}. \tag{2.18}$$

with Majorana mass of right-handed neutrino  $M_N = y^N v_s$  can be realized.<sup>17)</sup> This is so-called Type-I seesaw mechanism in a TeV scale, when coefficients  $y^\nu$  and  $y^N$  are assumed to be of order one. The masses of scalar and pseudo-scalar mostly from  $S$  are given by

$$m_{H_S}^2 = m_3^2 + 2\lambda_S v_s^2, \quad m_{A_S}^2 = 9\lambda v_s, \tag{2.19}$$

in the potential Eq. (2.16) without CP violation. For parameter region with  $v_s \gg 1$  TeV, both scalar and pseudo-scalar are heavier than other particles. After integrating out  $S$ , thanks to the  $Z_3$ -symmetry, the model ends up with an effectively neutrinophilic THDM with approximated  $Z_2$ -symmetry,  $\Phi \rightarrow \Phi, \Phi_\nu \rightarrow -\Phi_\nu$ . Comparing to the neutrinophilic THDM, the value of  $m_3^2$ , which is a soft  $Z_2$ -symmetry breaking term, is expected to be  $\kappa v_s$ .  $\lambda_5$  is induced by integrating out  $S$ , which is estimated as  $\mathcal{O}(\kappa^2/m_S^2) \sim 10^{-12}$ . Thus, the neutrinophilic THDM has an approximate  $Z_2$ -symmetry.

As for the neutrino mass induced from one-loop diagram <sup>\*)</sup>, UV theory induces small  $\lambda_5 \sim 10^{-12}$  due to  $Z_3$ -symmetry, so that radiative induced neutrino mass from one-loop diagram is estimated as  $\lambda_5 v^2 / (4\pi)^2 M \sim 10^{-4}$  eV. This can be negligible comparing to light neutrino mass which is induced from tree level Type-I seesaw mechanism. The tree level neutrino mass is

$$m_\nu^{tree} \sim \frac{y_\nu^2 v_\nu^2}{M} \sim \frac{y_\nu^2 \kappa^2 v^2}{v_s^2 M}, \quad (2.20)$$

where we input  $v_\nu \sim \frac{\kappa v}{v_s}$ . On the other hand, one-loop induced neutrino mass is estimated as

$$m_\nu^{loop} \sim \frac{\lambda_5 y_\nu^2 v^2}{16\pi^2 M} \sim \frac{y_\nu^2}{16\pi^2} \frac{\kappa^2 v^2}{M^2 M}. \quad (2.21)$$

Putting  $M \sim v_s$ ,

$$\frac{m_\nu^{loop}}{m_\nu^{tree}} \sim \frac{1}{16\pi^2}, \quad (2.22)$$

which shows loop induced neutrino mass is always smaller than tree level mass if UV theory is the model of Ref.<sup>17)</sup>

### 2.3. Supersymmetric extension of neutrinophilic Higgs doublet model

Now let us show the supersymmetric extension of the neutrinophilic Higgs doublet model. The supersymmetric extension is straightforward by extending its Higgs sector to be a four Higgs doublet model. The superpotential is given by

$$\begin{aligned} \mathcal{W} = & y^u \bar{Q}^L H_u U_R + y^d \bar{Q}^L H_d D_R + y^l \bar{L} H_d E_R + y^\nu \bar{L} H_\nu N + M N^2 \\ & + \mu H_u H_d + \mu' H_\nu H_{\nu'} + \rho H_u H_{\nu'} + \rho' H_\nu H_d, \end{aligned} \quad (2.23)$$

where  $H_u$  ( $H_d$ ) is Higgs doublet which gives mass of up- (down-) sector.  $H_\nu$  gives neutrino Dirac mass and  $H_{\nu'}$  does not contribute to fermion masses. For the  $Z_2$ -parity,  $H_u, H_d$  are even, while  $H_\nu, H_{\nu'}$  are odd. The  $Z_2$ -partity is softly broken by the  $\rho$  and  $\rho'$ . We assume that  $|\mu|, |\mu'| \gg |\rho|, |\rho'|$ , and SUSY breaking soft squared masses can trigger suitable electro-weak symmetry breaking. The Higgs potential is given by

$$\begin{aligned} V = & (|\mu|^2 + |\rho|^2) H_u^\dagger H_u + (|\mu|^2 + |\rho'|^2) H_d^\dagger H_d + (|\mu'|^2 + |\rho|^2) H_\nu^\dagger H_\nu + (|\mu'|^2 + |\rho|^2) H_{\nu'}^\dagger H_{\nu'} \\ & + \frac{g_1^2}{2} \left( H_u^\dagger \frac{1}{2} H_u - H_d^\dagger \frac{1}{2} H_d + H_\nu^\dagger \frac{1}{2} H_\nu - H_{\nu'}^\dagger \frac{1}{2} H_{\nu'} \right)^2 \\ & + \sum_a \frac{g_2^2}{2} \left( H_u^\dagger \frac{\tau^a}{2} H_u + H_d^\dagger \frac{\tau^a}{2} H_d + H_\nu^\dagger \frac{\tau^a}{2} H_\nu + H_{\nu'}^\dagger \frac{\tau^a}{2} H_{\nu'} \right)^2 \\ & + m_{H_u}^2 H_u^\dagger H_u + m_{H_d}^2 H_d^\dagger H_d + m_{H_\nu}^2 H_\nu^\dagger H_\nu + m_{H_{\nu'}}^2 H_{\nu'}^\dagger H_{\nu'} \end{aligned}$$

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<sup>\*)</sup> We would like to thank J. Kubo and H. Sugiyama for letting us notice this topic.

$$\begin{aligned}
& +B\mu H_u \cdot H_d + B'\mu' H_\nu \cdot H_{\nu'} + \hat{B}\rho H_u \cdot H_{\nu'} + \hat{B}'\rho' H_\nu \cdot H_d \\
& +\mu^* \rho H_d^\dagger H_{\nu'} + \mu^* \rho' H_u^\dagger H_\nu + \mu'^* \rho' H_{\nu'}^\dagger H_d + \mu'^* \rho H_\nu^\dagger H_u + h.c.,
\end{aligned} \tag{2.24}$$

where  $\tau^a$  and dot represent a generator of  $SU(2)$  and its anti-symmetric product respectively. We assume  $\text{Max.}[|\hat{B}\rho|, |\hat{B}'\rho'|, |\mu\rho|, |\mu'\rho|, |\mu\rho'|, |\mu'\rho'|] \sim \mathcal{O}(10) \text{ GeV}^2$ , which triggers VEV hierarchy between the SM Higgs doublet and neutrinophilic Higgs doublets. Notice that quarks and charged lepton have small non-holomorphic Yukawa couplings with  $H_\nu$ , through one-loop diagrams associated with small mass parameters of  $\hat{B}\rho, \hat{B}'\rho', \mu\rho, \mu'\rho, \mu\rho', \mu'\rho'$ . This situation is quite different from non-SUSY model, where these couplings are extremely suppressed by factor of  $v_\nu/v$ . As for the gauge coupling unification, we must introduce extra particles, but anyhow, the supersymmetric extension of neutrinophilic Higgs doublet model can be easily constructed as shown above.

### §3. Leptogenesis

#### 3.1. A brief overview of thermal leptogenesis

In the seesaw model, the smallness of the neutrino masses can be naturally explained by the small mixing between left-handed neutrinos and heavy right-handed Majorana neutrinos  $N_i$ . The basic part of the Lagrangian in the SM with right-handed neutrinos is described as

$$\mathcal{L}_N^{\text{SM}} = -y_{ij}^\nu \bar{l}_{L,i} \Phi N_j - \frac{1}{2} \sum_i M_i \bar{N}_i^c N_i + h.c., \tag{3.1}$$

where  $i, j = 1, 2, 3$  denote the generation indices,  $h$  is the Yukawa coupling,  $l_L$  and  $\Phi$  are the lepton and the Higgs doublets, respectively, and  $M_i$  is the lepton-number-violating mass term of the right-handed neutrino  $N_i$  (we are working on the basis of the right-handed neutrino mass eigenstates). With this Yukawa couplings, the mass of left-handed neutrino is expressed by the well-known formula

$$m_{ij} = \sum_k \frac{y_{ik}^\nu y_{kj}^{\nu T}}{M_k}. \tag{3.2}$$

The decay rate of the lightest right-handed neutrino is given by

$$\Gamma_{N_1} = \sum_j \frac{y_{1j}^{\nu\dagger} y_{j1}^\nu}{8\pi} M_1 = \frac{(y^{\nu\dagger} y^\nu)_{11}}{8\pi} M_1. \tag{3.3}$$

Comparing to the Friedmann equation for a spatially flat spacetime

$$H^2 = \frac{1}{3M_P^2} \rho, \tag{3.4}$$



with the energy density of the radiation

$$\rho = \frac{\pi^2}{30} g_* T^4, \quad (3.5)$$

where  $g_*$  denotes the effective degrees of freedom of relativistic particles and  $M_P \simeq 2.4 \times 10^{18}$  GeV is the reduced Planck mass, the condition of the out of equilibrium decay  $\Gamma_{N_1} < H|_{T=M_1}$  is rewritten as

$$\tilde{m}_1 \equiv (y^{\nu\dagger} y^\nu)_{11} \frac{v^2}{M_1} < \frac{8\pi v^2}{M_1^2} H|_{T=M_1} \equiv m_* \quad (3.6)$$

with  $m_* \simeq 1 \times 10^{-3}$  eV and  $v = 174$  GeV.

In the case of the hierarchical mass spectrum for right-handed neutrinos, the lepton asymmetry in the Universe is generated dominantly by CP-violating out of equilibrium decay of the lightest heavy neutrino,  $N_1 \rightarrow l_L \Phi^*$  and  $N_1 \rightarrow \bar{l}_L \Phi$ . Then, its CP asymmetry is given by<sup>23)</sup>

$$\begin{aligned} \varepsilon &\equiv \frac{\Gamma(N_1 \rightarrow \Phi + \bar{l}_j) - \Gamma(N_1 \rightarrow \Phi^* + l_j)}{\Gamma(N_1 \rightarrow \Phi + \bar{l}_j) + \Gamma(N_1 \rightarrow \Phi^* + l_j)} \\ &\simeq -\frac{3}{8\pi} \frac{1}{(y^\nu y^{\nu\dagger})_{11}} \sum_{i=2,3} \text{Im}(y^\nu y^{\nu\dagger})_{1i}^2 \frac{M_1}{M_i}, \quad \text{for } M_i \gg M_1. \end{aligned} \quad (3.7)$$

Through the relations of the seesaw mechanism, this can be roughly estimated as

$$\varepsilon \simeq \frac{3}{8\pi} \frac{M_1 m_3}{v^2} \sin \delta \simeq 10^{-6} \left( \frac{M_1}{10^{10} \text{GeV}} \right) \left( \frac{m_3}{0.05 \text{eV}} \right) \sin \delta, \quad (3.8)$$

where  $m_3$  is the heaviest light neutrino mass normalized by 0.05 eV which is a preferred to account for atmospheric neutrino oscillation data.<sup>24)</sup>

Using the above  $\varepsilon$ , the resultant baryon asymmetry generated via thermal leptogenesis is expressed as

$$\frac{n_b}{s} \simeq C \kappa \frac{\varepsilon}{g_*}, \quad (3.9)$$

where  $g_*|_{T=M_1} \sim 100$ , the so-called dilution (or efficiency) factor  $\kappa \leq \mathcal{O}(0.1)$  denotes the dilution by wash out processes, the coefficient

$$C = \frac{8N_f + 4N_H}{22N_f + 13N_H}, \quad (3.10)$$

with  $N_f$  and  $N_H$  being the number of fermion generation and Higgs doublet<sup>25)</sup> is the factor of the conversion from lepton to baryon asymmetry by the sphaleron.<sup>26)</sup> In order to obtain the observed baryon asymmetry in our Universe  $n_b/s \simeq 10^{-10}$ ,<sup>27)</sup> the inequality  $\varepsilon \gtrsim 10^{-7}$  is required. This can be rewritten as  $M_1 \gtrsim 10^9$  GeV, which is the so-called Davidson-Ibarra bound for models with hierarchical right-handed neutrino mass spectrum.<sup>4), 5)</sup>

### 3.2. leptogenesis in neutrinophilic THDM

Now we consider leptogenesis in the neutrinophilic THDM with the extra Higgs doublet  $\Phi_\nu$  described in Sec. 2.1. The relevant interaction part of Lagrangian Eq. (2.1) is expressed as

$$\mathcal{L}_N = -y_{ij}^\nu \bar{l}_{L,i} \Phi_\nu N_j - \frac{1}{2} \sum_i M_i \bar{N}_i^c N_i + h.c.. \quad (3.11)$$

The usual Higgs doublet  $\Phi$  in Eq. (3.1) is replaced by new Higgs doublet  $\Phi_\nu$ . Again, we are working on the basis of the right-handed neutrino mass eigenstates. Then, with these Yukawa couplings, the mass of left-handed neutrino is given by

$$m_{ij} = \sum_k \frac{y_{ik}^\nu v_\nu y_{kj}^{\nu T} v_\nu}{M_k}. \quad (3.12)$$

Thus, for a smaller VEV of  $v_\nu$ , a larger  $y^\nu$  is required.

The Boltzmann equation for the lightest right-handed neutrino  $N_1$ , which is denoted by  $N$  here, is given by

$$\begin{aligned} \dot{n}_N + 3Hn_N = & -\gamma(N \rightarrow L\Phi_\nu) - \gamma(N \rightarrow \bar{L}\Phi_\nu^*) \quad : \text{decay} \\ & +\gamma(L\Phi_\nu \rightarrow N) + \gamma(\bar{L}\Phi_\nu^* \rightarrow N) \quad : \text{inverse decay} \\ & -\gamma(NL \rightarrow A\Phi_\nu) - \gamma(N\Phi_\nu \rightarrow LA) - \gamma(N\bar{L} \rightarrow A\Phi_\nu^*) - \gamma(N\Phi_\nu^* \rightarrow \bar{L}A) \\ & + \text{inverse processes} \quad : \text{s-channel scattering} \\ & -\gamma(NL \rightarrow A\Phi_\nu) - \gamma(N\Phi_\nu \rightarrow LA) - \gamma(NA \rightarrow L\Phi_\nu) \\ & -\gamma(N\bar{L} \rightarrow A\Phi_\nu^*) - \gamma(N\Phi_\nu^* \rightarrow \bar{L}A) - \gamma(NA \rightarrow \bar{L}\Phi_\nu^*) \\ & + \text{inverse processes} \quad : \text{t-channel scattering} \\ & -\gamma(NN \rightarrow \text{Final}) + \gamma(\text{Final} \rightarrow NN) : \text{annihilation} \\ = & -\Gamma_D(n_N - n_N^{eq}) - \Gamma_{scat}(n_N - n_N^{eq}) - \langle \sigma v(\rightarrow \Phi, \Phi_\nu) \rangle (n_N^2 - n_N^{eq2}) \end{aligned} \quad (3.13)$$

where  $\Phi, \Phi_\nu$  and  $A$  denote the Higgs bosons, the neutrinophilic Higgs bosons and gauge bosons, respectively. Notice that usual  $\Delta L = 1$  lepton number violating scattering processes involving top quark is absent in this model, because  $\Phi_\nu$  has neutrino Yukawa couplings. Although the annihilation processes ( $NN \rightarrow \text{Final}$ ) is noted in Eq. (3.13), in practice, this is not relevant because the coupling  $y_{i1}^\nu$  is so small, as will be shown later, to satisfy the out of equilibrium decay condition.

The Boltzmann equation for the lepton asymmetry  $L \equiv l - \bar{l}$  is given by

$$\dot{n}_L + 3Hn_L$$

$$\begin{aligned}
&= \gamma(N \rightarrow l\Phi_\nu) - \gamma(\bar{N} \rightarrow \bar{l}\Phi_\nu^*) \\
&\quad - \{\gamma(l\Phi_\nu \rightarrow N) - \gamma(\bar{l}\Phi_\nu^* \rightarrow \bar{N})\} \quad : \text{decay and inverse decay} \\
&\quad - \gamma(lA \rightarrow N\Phi_\nu) + \gamma(\bar{l}A \rightarrow \bar{N}\Phi_\nu^*) - \gamma(Nl \rightarrow A\Phi_\nu) \\
&\quad + \gamma(\bar{N}\bar{l} \rightarrow A\Phi_\nu^*) \quad : \text{s-channel } \Delta L = 1 \text{ scattering} \\
&\quad - \gamma(Nl \rightarrow A\Phi_\nu) + \gamma(\bar{N}\bar{l} \rightarrow A\Phi_\nu^*) - \gamma(lA \rightarrow N\Phi_\nu) + \gamma(\bar{l}A \rightarrow \bar{N}\Phi_\nu^*) \\
&\quad - \gamma(l\Phi_\nu \rightarrow NA) + \gamma(\bar{l}\Phi_\nu^* \rightarrow \bar{N}A) \quad : \text{t-channel } \Delta L = 1 \text{ scattering} \\
&\quad + \gamma(\bar{l}\bar{l} \rightarrow \Phi_\nu^*\Phi_\nu^*) - \gamma(ll \rightarrow \Phi_\nu\Phi_\nu) \\
&\quad + 2\{\gamma(\bar{l}\Phi_\nu^* \rightarrow l\Phi_\nu) - \gamma(l\Phi_\nu \rightarrow \bar{l}\Phi_\nu^*)\} \quad : \text{t and s-channel } \Delta L = 2 \text{ scattering} \\
&= \varepsilon\Gamma_D(n_N - n_N^{eq}) - \Gamma_W n_L
\end{aligned} \tag{3.14}$$

where

$$\Gamma_W = \frac{1}{2} \frac{n_N^{eq}}{n_\gamma} \Gamma_N + \frac{n_N}{n_N^{eq}} \Gamma_{\Delta L=1,t} + 2\Gamma_{\Delta L=1,s} + 2\Gamma_{\Delta L=2} \tag{3.15}$$

is the wash-out rate.

The condition of the out of equilibrium decay is given as

$$\tilde{m}_1 \equiv (y^{\nu\dagger}y^\nu)_{11} \frac{v_\nu^2}{M_1} < \frac{8\pi v_\nu^2}{M_1^2} H|_{T=M_1} \equiv m_* \left(\frac{v_\nu}{v}\right)^2 \tag{3.16}$$

Notice that for  $v_\nu \ll v$  the upper bound on  $\tilde{m}_1$  becomes more stringent, which implies that the lightest left-handed neutrino mass is almost vanishing  $m_1 \simeq 0$ . Alternatively the condition can be expressed as

$$(y^{\nu\dagger}y^\nu)_{11} < 8\pi \sqrt{\frac{\pi^2 g_*}{90}} \frac{M_1}{M_P}. \tag{3.17}$$

Hence, for the TeV scale  $M_1$ , the value of  $(y^{\nu\dagger}y^\nu)_{11}$  must be very small, which can be realized by taking all  $y_{i1}^\nu$  to be small. Under such neutrino Yukawa couplings  $y_{i1}^\nu \ll y_{i2}^\nu, y_{i3}^\nu$  and hierarchical right-handed neutrino mass spectrum, the CP asymmetry,

$$\begin{aligned}
\varepsilon &\simeq -\frac{3}{8\pi} \frac{1}{(y^{\nu\dagger}y^\nu)_{11}} \left( \text{Im}(y^{\nu\dagger}y^\nu)_{12}^2 \frac{M_1}{M_2} + \text{Im}(y^{\nu\dagger}y^\nu)_{13}^2 \frac{M_1}{M_3} \right) \\
&\simeq -\frac{3}{8\pi} \frac{m_\nu M_1}{v_\nu^2} \sin\theta \\
&\simeq -\frac{3}{16\pi} 10^{-6} \left( \frac{0.1\text{GeV}}{v_\nu} \right)^2 \left( \frac{M_1}{100\text{GeV}} \right) \left( \frac{m_\nu}{0.05\text{eV}} \right) \sin\theta,
\end{aligned} \tag{3.18}$$

is significantly enhanced due to large Yukawa couplings  $y_{2i}^\nu$  and  $y_{3i}^\nu$  as well as the tiny Higgs VEV  $v_\nu$ . The thermal averaged interaction rate of  $\Delta L = 2$  scatterings is expressed as

$$\Gamma^{(\Delta L=2)} = \frac{1}{n_\gamma} \frac{T}{32\pi(2\pi)^4} \int ds \sqrt{s} K_1 \left( \frac{\sqrt{s}}{T} \right) \int \frac{d\cos\theta}{2} \sum |\mathcal{M}|^2 \tag{3.19}$$

with

$$\sum |\overline{\mathcal{M}}|^2 = 2|\overline{\mathcal{M}}_t|^2 + 2|\overline{\mathcal{M}}_s|^2 \simeq \sum_{j,(\alpha,\beta)} 2|y_{\alpha j}^\nu y_{\beta j}^{\nu\dagger}| \frac{s}{M_{N_j}^2}, \quad \text{for } s \ll M_j^2. \quad (3.20)$$

The decoupling condition

$$\Gamma^{(\Delta L=2)} < \sqrt{\frac{\pi^2 g_*}{90}} \frac{T^2}{M_P}, \quad (3.21)$$

for  $T < M_1$  is rewritten as

$$\sum_i \left( \sum_j \frac{y_{ij}^\nu y_{ji}^{\nu\dagger} v_\nu^2}{M_j} \right)^2 < 32\pi^3 \zeta(3) \sqrt{\frac{\pi^2 g_*}{90}} \frac{v_\nu^4}{T M_P}. \quad (3.22)$$

For lower  $v_\nu$ ,  $\Delta L = 2$  wash out processes are more significant. Inequality (3.22) gives the lower bound on  $v_\nu$  in order to avoid too strong wash out.

We here summarize all conditions for successful thermal leptogenesis, and the result is presented in Fig. 1. The horizontal axis is the VEV of neutrino Higgs  $v_\nu$  and the vertical axis is the mass of the lightest right-handed neutrino,  $M_1$ . In the red brown region, the lightest right-handed neutrino decay into Higgs boson  $H$  with assuming  $M_H = 100$  GeV, and lepton is kinematically not allowed. In turquoise region corresponds to inequality (3.22),  $\Delta L = 2$  wash out effect is too strong. The red and green line is contour of the CP asymmetry of  $\varepsilon = 10^{-6}$  and  $10^{-7}$ , respectively, with the lightest right-handed neutrino decay in hierarchical right-handed neutrino mass spectrum. Thus, in the parameter region above the line of  $\varepsilon = 10^{-7}$ , thermal leptogenesis easily works even with hierarchical masses of right-handed neutrinos. For the region below the line of  $\varepsilon = 10^{-7}$ , the resonant leptogenesis mechanism,<sup>10)</sup> where CP asymmetry is enhanced resonantly by degenerate right-handed neutrino masses, may work. Here we stress that, for  $v_\nu \ll 100$  GeV, the required degree of mass degeneracy is considerably milder than that for the original resonant leptogenesis.

### 3.3. Constraints on an UV theory

Let us suppose that neutrinophilic THDM is derived from a model reviewed in Sec. 2.2 by integrated out a singlet field  $S$ . If  $S$  is relatively light, thermal leptogenesis discussed above could be affected. That is the annihilation processes of  $N_1$  which has been justifiably ignored in Eq. (3.13). However, the annihilation could take place more efficiently via S-channel  $S$  scalar exchange processes in the UV theory.<sup>17)</sup>

For example, the annihilation  $N_1 N_1 \rightarrow \Phi_\nu \Phi_\nu^*$  with the amplitude

$$|\overline{\mathcal{M}}|^2 = \left| \frac{y_1^\nu \lambda_{\Phi_\nu} v_s}{s - M_{H_S}^2 - i M_{H_S} \Gamma_{H_S}} \right|^2 \frac{s - 4M_1^2}{4}, \quad (3.23)$$

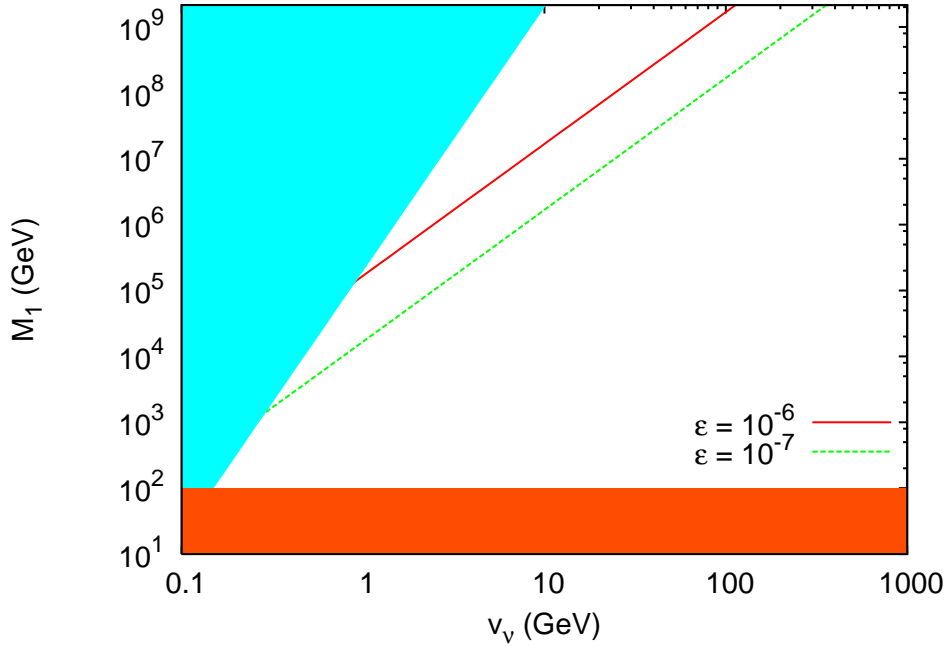


Fig. 1. Available region for leptogenesis. The horizontal axis is the VEV of neutrino Higgs  $v_\nu$  and the vertical axis is the mass of the lightest right-handed neutrino mass  $M_1$ . In the red brown region, the lightest right-handed neutrino decay into Higgs boson  $\Phi_\nu$  and lepton is kinematically forbidden. In turquoise region,  $\Delta L = 2$  wash out effect is too strong. The red and green line is contour of the CP asymmetry of  $\epsilon = 10^{-6}$  and  $10^{-7}$ , respectively, with the lightest right-handed neutrino decay in hierarchical right-handed neutrino mass spectrum.

would not be in equilibrium, if

$$\lambda_{\Phi_\nu} \lesssim 40 \frac{M_{H_S}^2}{M_1^{3/2} M_P^{1/2}} \simeq 0.1 \left( \frac{10^5 \text{ GeV}}{M_1} \right)^{3/2} \left( \frac{M_{H_S}}{10^7 \text{ GeV}} \right)^2, \quad (3.24)$$

for  $M_S \gg T > M_1$  is satisfied. Here,  $\Gamma_{H_S}$  denotes the decay width of  $H_S$ . Constraints on other parameters such as  $\lambda_\Phi$  and  $\kappa$  can be similarly obtained.

#### §4. Supersymmetric case: Reconciling to thermal leptogenesis, gravitino problem and neutralino dark matter

As we have shown in Sec. 2.3, it is possible to construct a supersymmetric model with  $\Phi_\nu$ . A discrete symmetry, called “R-parity”, is imposed in many supersymmetric models in order to prohibit rapid proton decay. Another advantage of the conserved R-parity is that it guarantees the absolute stability of the LSP, which becomes a dark matter candidate. In large parameter space of supergravity model with gravity mediated SUSY breaking, gravitino has the mass of  $\mathcal{O}(100)$  GeV and decays into LSP (presumably the lightest neutralino) at

late time after BBN. Then, decay products may affect the abundances of light elements produced during BBN. This is so-called “gravitino problem”.<sup>6)</sup> To avoid this problem, the upper bound on the reheating temperature after inflation

$$T_R < 10^6 - 10^7 \text{ GeV}, \quad (4.1)$$

has been derived as depending on gravitino mass.<sup>7)</sup> By comparing Eq. (4.1) with the CP violation in supersymmetric models with hierarchical right-handed neutrino masses, which is about four times larger than that in non-supersymmetric model<sup>28)</sup> as,

$$\varepsilon \simeq -\frac{3}{2\pi} \frac{1}{(y^\nu y^{\nu\dagger})_{11}} \sum_{i=2,3} \text{Im}(y^\nu y^{\nu\dagger})_{1i}^2 \frac{M_1}{M_i}, \quad (4.2)$$

it has been regarded that thermal leptogenesis through the decay of heavy right-handed neutrinos hardly work because of gravitino problem.

As we have shown in the previous section, a sufficient CP violation  $\varepsilon = \mathcal{O}(10^{-6})$  can be realized for  $v_\nu = \mathcal{O}(1)$  GeV in the hierarchical right-handed neutrino masses with  $M_1$  of  $\mathcal{O}(10^5 - 10^6)$  GeV. This implies that the reheating temperature after inflation  $T_R$  of  $\mathcal{O}(10^6)$  GeV is high enough in order to produce right-handed neutrinos by thermal scatterings. Thus, it is remarkable that SUSY neutrinophilic model with  $v_\nu = \mathcal{O}(1)$  GeV can realize thermal leptogenesis in gravity mediated SUSY breaking with unstable gravitino. In this setup, the lightest neutralino could be LSP and dark matter with the standard thermal freeze out scenario.

## §5. Conclusion

We have examined the possibility of thermal leptogenesis in neutrinophilic Higgs doublet models, whose tiny VEV gives neutrino Dirac mass term. Thanks to the tiny VEV of the neutrinophilic Higgs field, neutrino Yukawa couplings are not necessarily small, instead, they tend to be large, and the CP asymmetry in the lightest right-handed neutrino decay is significantly enlarged. Although the  $\Delta L = 2$  wash out effect also could be enhanced simultaneously, we have found the available parameter region where its wash out effect is avoided as keeping the CP asymmetry large enough. In addition, in a supersymmetric neutrinophilic Higgs doublet model, we have pointed out that thermal leptogenesis in gravity mediated SUSY breaking works well without confronting gravitino problem. Where the lightest neutralino could be LSP and dark matter with the standard thermal freeze out scenario.

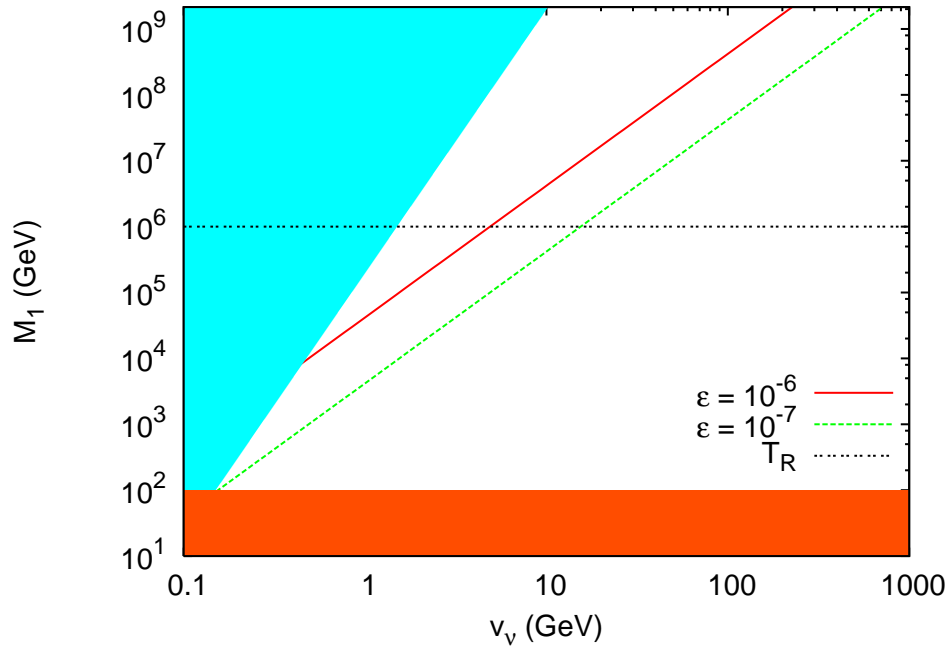


Fig. 2. The same as Fig. 1 but with Eq. (4.2). The additional horizontal black dashed line represents a reference value of the upper bound on reheating temperature after inflation  $T_R$  of  $10^6$  GeV from gravitino overproduction.

### Acknowledgements

We would like to thank M. Hirotsu for collaboration in the early stage of this work. We are grateful to S. Matsumoto, S. Kanemura and K. Tsumura for useful and helpful discussions. This work is partially supported by Scientific Grant by Ministry of Education and Science, Nos. 20540272, 22011005, 20039006, 20025004 (N.H.), and the scientific research grants from Hokkai-Gakuen (O.S.).

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